

The Diphoton Excess, Low Energy Theorem and the 331 Model

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Abstract

We interpret the diphoton anomaly as a heavy scalar H_3 in the so-called 331 model. The scalar is responsible for breaking the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge symmetry down to the standard model electroweak gauge group. It mainly couples to the standard model gluons and photons through quantum loops involving heavy quarks and leptons. Those quarks and leptons, in together with the SM quarks and leptons, form the fundamental representation of the 331 model. We use low energy theorem to calculate effective couplings of $H_3 gg$, $H_3 \gamma\gamma$, $H_3 ZZ$, $H_3 WW$ and $H_3 Z\gamma$. The analytical results can be applied to new physics models satisfying the low energy theorem. We show that the heavy quark and lepton contribution cannot produce enough diphoton pairs. It is crucial to include the contribution of charged scalars to explain the diphoton excess. The extra neutral Z' boson could also explain the 2 TeV diboson excess observed at the LHC Run-I.

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I. INTRODUCTION

A diphoton resonance around 750 GeV was reported by the ATLAS and CMS collaborations [1, 2] based on 3.2 and 2.6 fb⁻¹ of data collected at 13 TeV LHC. ATLAS observed a local significance of 3.9 σ corresponding to a cross section of about 10 ± 3 fb, while CMS found 2.6 σ with a smaller cross section of 6 ± 3 fb. And the large excess around 750 GeV is still confirmed by a recent study from ATLAS and CMS [3, 4]. The resonance is likely to be a new scalar in new physics (NP) model beyond the Standard Model (SM) of particle physics. It has drawn a lot of interests in the field [5–96]. Also, an excess of diboson events around 2 TeV was reported in LHC Run-I data [97–99]. That has been explained in terms of extra gauge bosons such as additional Z' or W' bosons in several NP models with extended gauge structures. New physics models that could explain both anomalies need to have two new ingredients, extra heavy scalars and gauge bosons. The simplest case is the new scalar is singlet under the SM gauge group which requires the model has vector-like fermions in order to couple to the diphoton at the loop-level. Many papers add the singlet scalar and vector-like fermions by hand, while in the so-called 331 model with the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge group [100, 101] these additional particles naturally exist. There are many versions of the 331 model [102–110] which in general share nice features as follows: i) the anomaly cancellation and the QCD asymptotic freedom require three generation fermions; ii) the Peccei-Quinn (PQ) symmetry [111, 112] which can solve the strong CP problem is a natural result of gauge invariance in the 331 model [102, 103], etc. Refs. [70, 72] studied the diphoton excess in the framework of the 331 gauge symmetry in a specific version of the 331 model while in our paper we consider four possible versions of this model. Also, the trinification model of $SU(3)_L \otimes SU(3)_R \otimes SU(3)_C$, which contains the 331 model as subgroup, has been considered recently in Ref. [74] to address on the diphoton anomaly.

The 331 model consists of very rich particle spectra. For instance, there will be five new gauge bosons, three new heavy quarks, three new leptons, and six new scalars. That has very rich collider phenomenology at the LHC. There are many versions of the 331 model which can be characterized by a parameter called β . Models with different β have new particles with different electric charges. The 331 model has been studied in details by Buras *et al* in Ref. [113]. In this work we follow the notation in Ref. [113] and consider different versions of the 331 model with $\beta = \pm\sqrt{3}$ and $\pm 1/\sqrt{3}$ to explain the 750 GeV diphoton and

2 TeV diboson excesses.

The paper is organized as follows. In Sec. II, we briefly review the 331 model and introduce the main ingredients of the 331 model needed to explain the diphoton and diboson excesses at the LHC. In Sec. III we explain how the 331 model could explain the diphoton and diboson signal. Finally, we conclude in Sec. IV.

II. THE MODEL

In the 331 model considered in this work, the right-handed fermion fields are treated as singlets of the $SU(3)_L$ group. While the left-handed quark fields of the first two generations are chosen to be in the triplet representation of the $SU(3)_L$ group and the left-handed quarks of the third generation are in the anti-triplet representation of $SU(3)_L$. Three new quarks (D, S, T) and three new leptons (E_e, E_μ, E_τ) are added to form the triplet and anti-triplet. The quark fields are

$$q_{1L} = \begin{pmatrix} u \\ d \\ D \end{pmatrix}_L, \quad q_{2L} = \begin{pmatrix} c \\ s \\ S \end{pmatrix}_L, \quad q_{3L} = \begin{pmatrix} b \\ -t \\ T \end{pmatrix}_L. \quad (1)$$

Note that the t and b assignment is different from the SM as a result of requiring q_{3L} being an anti-triplet. All the left-handed lepton fields of the three generations are all treated as anti-triplets to guarantee the gauge anomaly cancellation, which requires equal numbers of triplets and anti-triplets. The lepton fields are given by

$$l_{1L} = \begin{pmatrix} e \\ -\nu_e \\ E_e \end{pmatrix}_L, \quad l_{2L} = \begin{pmatrix} \mu \\ -\nu_\mu \\ E_\mu \end{pmatrix}_L, \quad l_{3L} = \begin{pmatrix} \tau \\ -\nu_\tau \\ E_\tau \end{pmatrix}_L, \quad (2)$$

The extra minus signs in front of top quark and neutrinos are to generate the same Feynman vertices as in the SM.

The symmetry breaking pattern of the 331 model is

$$SU(3)_L \times U(1)_X \rightarrow SU(2)_L \times U(1)_Y \rightarrow U(1)_Q, \quad (3)$$

which is realized by introducing three scalar triplets ρ, η and χ

$$\rho = \begin{pmatrix} \rho^+ \\ \rho^0 \\ \rho^{-Q_V} \end{pmatrix}, \quad \eta = \begin{pmatrix} \eta^0 \\ \eta^- \\ \eta^{-Q_Y} \end{pmatrix}, \quad \chi = \begin{pmatrix} \chi^{Q_Y} \\ \chi^{Q_V} \\ \chi^0 \end{pmatrix}, \quad (4)$$

TABLE I. *Electric charges of new particles for different choices of β .*

particles	$Q(\beta)$	$\beta = -\frac{1}{\sqrt{3}}$	$\beta = \frac{1}{\sqrt{3}}$	$\beta = -\sqrt{3}$	$\beta = \sqrt{3}$
D, S	$\frac{1}{6} - \frac{\sqrt{3}\beta}{2}$	$\frac{2}{3}$	$-\frac{1}{3}$	$\frac{5}{3}$	$-\frac{4}{3}$
T	$\frac{1}{6} + \frac{\sqrt{3}\beta}{2}$	$-\frac{1}{3}$	$\frac{2}{3}$	$-\frac{4}{3}$	$\frac{5}{3}$
E	$-\frac{1}{2} + \frac{\sqrt{3}\beta}{2}$	-1	0	-2	1
V	$-\frac{1}{2} + \frac{\sqrt{3}\beta}{2}$	-1	0	-2	1
Y	$\frac{1}{2} + \frac{\sqrt{3}\beta}{2}$	0	1	-1	2
H_V	$-\frac{1}{2} + \frac{\sqrt{3}\beta}{2}$	-1	0	-2	1
H_Y	$\frac{1}{2} + \frac{\sqrt{3}\beta}{2}$	0	1	-1	2
H_W	1	1	1	1	1

where the charges of Q_V and Q_Y are to be determined below. The vacuum expectation values (vevs) of ρ , η and χ are chosen as

$$\langle \rho \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_1 \\ 0 \end{pmatrix} \quad \langle \eta \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_2 \\ 0 \\ 0 \end{pmatrix} \quad \langle \chi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ v_3 \end{pmatrix}. \quad (5)$$

It is necessary to mention that the scalar in higher dimensional representation is also possible. However it is impossible to write the Yukawa interaction of the sextet scalar and quarks in a gauge invariant way [113] which is essential for the production of the neutral scalar at LHC. Therefore we only consider the three triplets. At the first step of the symmetry breaking, χ is introduced to break $SU(3)_L \times U(1)_X$ to $SU(2)_L \times U(1)_Y$ at a very large scale v_3 , typically at TeV scale. At the second step, we use ρ and η to break $SU(2)_L \times U(1)_Y$ down to the residual $U(1)_Q$ electromagnetic symmetry at the weak scale, i.e. $v_1 \sim v_2 \sim m_W$. Thus, we have $v_3 \gg v_{1,2}$. After the first step of symmetry breaking, the hypercharge generator is a linear combination of X and T^8 as $Y = \beta T^8 + X$, where T^8 is diagonal generator of $SU(3)_L$ and X is the generator of $U(1)_X$. The electric charge generator is then given by

$$Q = T^3 + \beta T^8 + X. \quad (6)$$

Table I shows electric charges of new particles. There are exotic quarks (i.e. D, S, T) for $\beta = \pm\sqrt{3}$ whose electric charges are $4/3$ or $5/3$ [114]. For $\beta = 1/\sqrt{3}$, the heavy leptons E are neutral and can be identified with right-handed neutrino. By introducing additional sextet scalar [115] one can generate small neutrino mass through the seesaw mechanism.

The scalar potential of the three scalar triplets is

$$V_H = \mu_1^2(\rho^\dagger \rho) + \mu_2^2(\eta^\dagger \eta) + \mu_3^2(\chi^\dagger \chi) + \lambda_1(\rho^\dagger \rho)^2 + \lambda_2(\eta^\dagger \eta)^2 + \lambda_3(\chi^\dagger \chi)^2$$

$$\begin{aligned}
& + \lambda_{12}(\rho^\dagger \rho)(\eta^\dagger \eta) + \lambda_{13}(\rho^\dagger \rho)(\chi^\dagger \chi) + \lambda_{23}(\eta^\dagger \eta)(\chi^\dagger \chi) \\
& + \lambda'_{12}(\rho^\dagger \eta)(\eta^\dagger \rho) + \lambda'_{13}(\rho^\dagger \chi)(\chi^\dagger \rho) + \lambda'_{23}(\eta^\dagger \chi)(\chi^\dagger \eta) + \sqrt{2}f (\epsilon_{ijk} \rho^i \eta^j \chi^k + \text{h.c.}), \quad (7)
\end{aligned}$$

where $\lambda_{i(j)}^{(\prime)}$ denotes the dimensionless parameter while μ_i and the parameter f has a mass dimension. For simplicity, we assume f is proportional to v_3 such that the theory has no other scale, i.e. $f \equiv kv_3$ with $k \sim O(1)$. Three CP-even scalars and one CP-odd scalar (A) remain after symmetry breaking. The mass mixing matrix of the CP-even scalars is

$$M_H^2 = \begin{pmatrix} 2\lambda_1 v_1^2 + \frac{f v_2 v_3}{v_1} & -f v_3 + v_1 v_2 \lambda_{12} & -f v_2 + v_1 v_3 \lambda_{13} \\ -f v_3 + v_1 v_2 \lambda_{12} & 2\lambda_2 v_2^2 + \frac{f v_1 v_3}{v_2} & -f v_1 + v_2 v_3 \lambda_{23} \\ -f v_2 + v_1 v_3 \lambda_{13} & -f v_1 + v_2 v_3 \lambda_{23} & 2\lambda_3 v_3^2 + \frac{f v_1 v_2}{v_3} \end{pmatrix}. \quad (8)$$

In the limit of $v_3 \gg v_{1,2}$, the three eigenvalues are

$$M_h^2 = \lambda_1 v_1^2 + \lambda_2 v_2^2, \quad M_{H_2}^2 = \frac{(v_1^2 + v_2^2)}{v_1 v_2} f v_3, \quad M_{H_3}^2 = 2\lambda_3 v_3^2. \quad (9)$$

Here h , H_2 and H_3 are three mass eigenstates of the neutral scalars, which are related with the weak eigenstates by a rotation matrix

$$\begin{pmatrix} h \\ H_2 \\ H_3 \end{pmatrix} = \begin{pmatrix} \frac{v_1}{\sqrt{v_1^2 + v_2^2}} & -\frac{v_2}{\sqrt{v_1^2 + v_2^2}} & 0 \\ \frac{v_2}{\sqrt{v_1^2 + v_2^2}} & \frac{v_1}{\sqrt{v_1^2 + v_2^2}} & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} \xi_\rho \\ \xi_\eta \\ \xi_\chi \end{pmatrix}, \quad (10)$$

where ξ_ρ , ξ_η and ξ_χ are weak eigenstates of the CP-even scalars. One of the three CP-even scalars is identified as the 125 GeV Higgs boson (h) while the other two (H_2 and H_3) are much heavier. As a result of the two-step symmetry breaking, the H_3 scalar is mainly from the χ triplet while the h and H_2 scalars are from the ρ and η triplets. The mixing between the H_3 scalar and the Higgs boson h is suppressed such that decays of the H_3 scalar into the SM particles are negligible. On the other hand, the H_2 scalar has a sizable mixing with the Higgs boson and then could decay into a pair of SM fermions, W and Z , and Higgs bosons, if kinematically allowed.

The CP-odd scalar A mainly arises from ρ and η as the imaginary part of the neutral component in the χ triplet is eaten by the new Z' boson. The scalar A could decay into the SM quarks and leptons directly. There are three charged scalars after symmetry breaking, named as H_W , H_V and H_Y , whose electric charges are ± 1 , $\pm Q_V$ and $\pm Q_Y$, respectively. See Table I for the scalar charges for different values of β .

In addition to the W -boson and the Z -boson, four charged gauge bosons and one neutral gauge boson appear after symmetry breaking. The four charged gauge bosons carry the same charge as the charged scalars; see Table I. We thus name them as $Y^{\pm Q_Y}$ and $V^{\pm Q_V}$ for simplicity. All the new gauge bosons achieve their masses after the first step of symmetry breaking and their masses are naturally around v_3 . Below we summarize the mass spectra of the vector bosons and scalars in the limit of $v_3 \gg v_{1,2}$:

- Gauge bosons

$$\begin{aligned} M_W &= \frac{1}{2}g\sqrt{v_1^2 + v_2^2}, & M_Y = M_V &= \frac{1}{2}gv_3, \\ M_Z &= \frac{1}{c_W}M_W, & M_{Z'} &= \frac{2c_W}{\sqrt{3}[1 - (1 + \beta^2)s_W^2]}M_Y, \end{aligned} \quad (11)$$

where $c_W \equiv \cos \theta_W$ and $s_W \equiv \sin \theta_W$ with θ_W being the Weinberg mixing angle.

- Scalars

$$\begin{aligned} M_{H_W} &\approx M_{H_2} \approx M_A \approx v_3 \sqrt{\frac{k(v_1^2 + v_2^2)}{v_1 v_2}}, & M_{H_3} &= \sqrt{2\lambda_3}v_3, \\ M_{H_Y} &= v_3 \sqrt{k \frac{v_1}{v_2} + \frac{\lambda'_{23}}{2}}, & M_{H_V} &= v_3 \sqrt{k \frac{v_2}{v_1} + \frac{\lambda'_{13}}{2}}, \end{aligned} \quad (12)$$

where $k \equiv f/v_3 \sim \mathcal{O}(1)$ is understood. The degeneracy of A , H_2 and H_W is owing to the remaining $SU(2)$ symmetry after the first step of symmetry breaking.

The masses of new fermions (\mathcal{F}) originate from the following Yukawa Lagrangian,

$$-\mathcal{L}_{Yuk}^{\mathcal{F}} = y_{ik}^J \bar{q}_{iL} \chi J_{kR} + y_{33}^J \bar{q}_{3L} \chi^* J_{3R} + y_{mn}^E \bar{l}_{mL} \chi^* E_{nR} + h.c. \quad , \quad (13)$$

where i (k) run from 1 to 2 while j (m, n) run from 1 to 3. J_{kR} refers to the right-handed heavy quarks D and S , J_{3R} to T , and E_{nR} represents the right-handed heavy leptons $E_{e,\mu,\tau}$. Note that χ can only give mass to the third component of the triplet fermions (i.e. $E_e, E_\mu, E_\tau, D, S, T$), therefore, there is no direct couplings of χ to the SM particles. After the first step of the symmetry breaking, the new fermions obtain their masses from the vacuum expectation value of χ . For simplicity, we assume the mass matrix of the new fermions is diagonal, leading to the following masses,

$$M_D = \frac{y_{11}^J}{\sqrt{2}}v_3 \quad M_S = \frac{y_{22}^J}{\sqrt{2}}v_3 \quad M_T = \frac{y_{33}^J}{\sqrt{2}}v_3 \quad (14)$$

$$M_{E_e} = \frac{y_{11}^E}{\sqrt{2}}v_3 \quad M_{E_\mu} = \frac{y_{22}^E}{\sqrt{2}}v_3 \quad M_{E_\tau} = \frac{y_{33}^E}{\sqrt{2}}v_3 \quad (15)$$

For illustration we choose benchmark parameters

$$M_{E_e} = M_{E_\mu} = M_{E_\tau} = M_D = M_S = M_T = 800 \text{ GeV}. \quad (16)$$

III. DIPHOTON EXCESS

To explain the diphoton excess, the cross section $\sigma(pp \rightarrow X \rightarrow \gamma\gamma)$ is estimated roughly as $(8 \pm 5) \text{ fb}$ [1, 2] where X denotes the possible resonance state. In the 331 model the H_3 scalar arises mainly from the χ triplet which decouples from the electroweak sector in the SM. The H_3 scalar does not couple to the SM particles directly. However, it could be produced through the heavy quark loops and also decays via the loops of new heavy charged particles; see Fig. 1. Using the narrow width approximation (NWA), the cross section of $\sigma(pp \rightarrow H_3 \rightarrow \gamma\gamma)$ at the LHC with energy $\sqrt{s} = 13 \text{ TeV}$ could be parameterized as [10]

$$\sigma(pp \rightarrow H_3 \rightarrow \gamma\gamma) = \frac{\mathcal{C}_{gg}}{M_{H_3}\Gamma_S} \Gamma(H_3 \rightarrow gg) \Gamma(H_3 \rightarrow \gamma\gamma), \quad (17)$$

where $\mathcal{C}_{gg} = 2137$ denotes the dimensionless partonic integral of parton distribution functions, and M_{H_3} is the mass of the H_3 scalar resonance. The partial decay widths of H_3 are given by

$$\begin{aligned} \Gamma(H_3 \rightarrow gg) &= \frac{1}{8\pi} c_{gg}^2 M_{H_3}, \\ \Gamma(H_3 \rightarrow \gamma\gamma) &= \frac{1}{64\pi} (-c_{\gamma\gamma}^{\mathcal{F}} + c_{\gamma\gamma}^{\mathcal{V}} + c_{\gamma\gamma}^{\mathcal{S}})^2 M_{H_3}, \\ \Gamma(H_3 \rightarrow ZZ) &= \mathcal{P}\left(\frac{m_Z^2}{M_{H_3}^2}\right) \frac{1}{64\pi} (c_{ZZ}^{\mathcal{V}} - c_{ZZ}^{\mathcal{F}} + c_{ZZ}^{\mathcal{S}})^2 M_{H_3}, \\ \Gamma(H_3 \rightarrow WW) &= \mathcal{P}\left(\frac{m_W^2}{M_{H_3}^2}\right) \frac{1}{32\pi} (c_{WW}^{\mathcal{V}} - c_{WW}^{\mathcal{F}} + c_{WW}^{\mathcal{S}})^2 M_{H_3}, \\ \Gamma(H_3 \rightarrow Z\gamma) &= \frac{1}{32\pi} \left(1 - \frac{m_Z^2}{M_{H_3}^2}\right)^3 (c_{Z\gamma}^{\mathcal{V}} - c_{Z\gamma}^{\mathcal{F}} + c_{Z\gamma}^{\mathcal{S}})^2 M_{H_3}, \end{aligned} \quad (18)$$

where $\mathcal{P}(x) = \sqrt{1-4x}(1-4x+6x^2)$ is a factor correcting for the massive final states in the decay width. The corresponding coefficients denoted by c_{ij}^k can be derived from the Higgs low-energy theorem which will be shown in the following subsections. We have checked the expressions of the coefficients with those from Ref. [116–118]. Note that the total decay width Γ can be obtained by summing over the above partial decay widths.

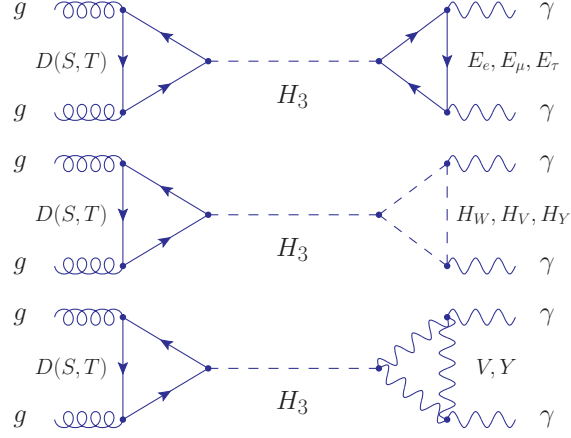


FIG. 1. The Feynman diagrams for $gg \rightarrow H \rightarrow \gamma\gamma$.

A. Loop-induced $H_3 \rightarrow gg$

The interaction of the H_3 scalar to the gluon pair through heavy quarks can be derived from the low energy theorem [119]. In the SM gluon fusion production is induced by the top quark loop [116],

$$c_{gg}^{\text{SM}} = \frac{\sqrt{2}\alpha_s}{3\pi} \frac{m_h}{v} \quad (19)$$

It is well-known that this coefficient is related to the top contribution to the gluon two-point function from the Higgs low-energy theorem. Next we consider the case when the loop contains a pair of vectorlike quarks (Q^c, Q) charged under the electroweak gauge group with the interaction

$$M_Q Q^c Q + y_Q H_3 Q^c Q, \quad (20)$$

then its contribution to the gluon two-point function is [118]

$$-\frac{1}{4} \left[1 - \frac{g_s^2}{16\pi^2} b_F^{(3)} \log \frac{M_Q^2(\mathcal{H})}{\mu^2} \right] G_{\mu\nu}^a G^{a\mu\nu}, \quad (21)$$

where $b_F^{(3)} = 2/3$ is the contribution to the one-loop beta function of QCD from a Dirac fermion and $M_Q(\mathcal{H}) = M_Q + y_Q \mathcal{H}$ is the mass of the new heavy fermion Q when turning on the scalar as a background field $H_3 \rightarrow \mathcal{H}$. The y_Q denotes the Yukawa coupling of the H_3 scalar to heavy quark Q . To obtain the scalar-gluon-gluon coupling, the Higgs low-energy theorem instructs us to expand Eq. 21 to the first order in \mathcal{H} [119]:

$$c_{gg} = \sum_{Q=D,S,T} \frac{\alpha_s}{3\pi} \frac{M_{H_3}}{M_Q} y_Q. \quad (22)$$

Strictly speaking, the low-energy theorem applies only when the mass of the particle in the loop is much larger than the scalar mass, $M_{H_3}^2/(4M_Q^2) \ll 1$, so that the loop diagram can be approximated by a dimension-five operator.

B. Loop induced $H_3 \rightarrow \gamma\gamma$

The decay of $H_3 \rightarrow \gamma\gamma$ can occur through a set of heavy fermions (\mathcal{F}), scalars (\mathcal{S}) and vector bosons (\mathcal{V}). First, the contribution of heavy quarks and leptons to the two-point function of the photon is

$$c_{\gamma\gamma}^{\mathcal{F}} = \sum_{\mathcal{F}} \frac{2\alpha N_c}{3\pi} Q_{\mathcal{F}}^2 \frac{M_{H_3}}{M_{\mathcal{F}}} y_{\mathcal{F}} \quad (23)$$

where \mathcal{F} denotes the heavy quarks (D, S, T) and leptons (E_e, E_μ, E_τ). Here, $Q_{\mathcal{F}}$ is the electric charge of those heavy fermions and $N_c = 3$ (1) for quarks (leptons) is understood. We note that there is a strong cancellation between new charged vector-bosons and new fermions, which leads to a small partial width of $H \rightarrow \gamma\gamma$.

In the 331 model the $H_3 \rightarrow \gamma\gamma$ can also be induced by a new set of charged vector-bosons V and Y . As the electric charges of V and Y depend on the choice of β , we introduce a general form of gauge boson couplings $\kappa g_{WW\gamma}^{\text{SM}}$ where $g_{WW\gamma}^{\text{SM}}$ is the $WW\gamma$ coupling in the SM. The contribution of charged vector bosons is simply

$$-\frac{1}{4} \left[1 - \frac{e^2 \kappa^2}{16\pi^2} 7 \log \frac{M_{\mathcal{V}}^2(\mathcal{H})}{\mu^2} \right] F_{\mu\nu} F^{\mu\nu}, \quad (24)$$

where $\mathcal{V} = V, Y$ and $F_{\mu\nu}$ is the field strength tensor of photon. Thus, we obtain the contribution of heavy new gauge bosons to the two-point function of photon is

$$c_{\gamma\gamma}^{\mathcal{V}} = \sum_{\mathcal{V}} \frac{7\alpha \kappa^2}{2\pi} \frac{M_{H_3}}{M_{\mathcal{V}}} y_{\mathcal{V}}. \quad (25)$$

In the SM the W contribution to the decay of $H \rightarrow \gamma\gamma$ dominates over the one from the top-quark loop owing to a large beta function coefficient “7” multiplying the W loop result.

Another contribution is from the charged scalar \mathcal{S} (H_W, H_V, H_Y) loops. For a general coupling of the photon to a pair of charged Higgs bosons,

$$g_{\gamma\mathcal{S}\mathcal{S}} = \lambda (p_1 - p_2)^\mu, \quad (26)$$

the loop contribution reads as

$$-\frac{1}{4} \left[1 + \frac{\lambda^2}{48\pi^2} \log \frac{M_S^2(\mathcal{H})}{\mu^2} \right] F_{\mu\nu} F^{\mu\nu}. \quad (27)$$

In the limit of $v_3 \gg v_{1,2}$, the charged scalar masses are proportional to v_3 , which leads to

$$c_{\gamma\gamma}^{\mathcal{S}} = - \sum_S \frac{\lambda^2}{24\pi^2} \frac{M_{H_3}}{M_S} y_S. \quad (28)$$

In the 331 model, the mass of charged scalar (H_W, H_V, H_Y) also depends on other parameters of the Higgs potential. The corresponding coefficient is then given by

$$c_{\gamma\gamma}^{\mathcal{S}}(331) = - \sum_S \frac{\lambda^2}{48\pi^2} \frac{\lambda_{\text{eff}} v_3}{M_S^2} M_{H_3}, \quad (29)$$

where λ_{eff} denotes the effective coupling of H_3 to charged scalars and M_S denotes the mass of charged scalar.

C. Loop induced $H_3 \rightarrow ZZ/WW$

For simplicity, we consider general interactions as follows:

$$\begin{aligned} g_{Z\nu\nu} &= \kappa' g_{ZWW}^{\text{SM}}, & g_{Z\mathcal{F}\mathcal{F}} &= g\gamma^\mu (g_L P_L + g_R P_R), \\ g_{ZSS} &= \lambda' (p_1 - p_2)^\mu, & g_{ZS\nu} &= \eta M_\nu g^{\mu\nu}. \end{aligned} \quad (30)$$

The contribution of the vector boson loop to the two-point function of the Z -boson is

$$-\frac{1}{4} \left[1 - \frac{e^2 \kappa'^2}{16\pi^2 s_W^2 c_W^2} 7c_W^4 \log \frac{M_V^2(\mathcal{H})}{\mu^2} \right] Z_{\mu\nu} Z^{\mu\nu}. \quad (31)$$

We expand Eq. 31 to the first order in H_3 to obtain the effective coefficient

$$c_{ZZ}^{\mathcal{V}} = \sum_{\mathcal{V}} \frac{7\alpha \kappa'^2}{2\pi} \frac{c_W^2}{s_W^2} \frac{M_{H_3}}{M_{\mathcal{V}}} y_{\mathcal{V}}. \quad (32)$$

The fermion loop contribution is

$$-\frac{1}{4} \left[1 - \frac{e^2 N_c}{24\pi^2 s_W^2} (g_L^2 + g_R^2) \log \frac{M_{\mathcal{F}}^2(\mathcal{H})}{\mu^2} \right] Z_{\mu\nu} Z^{\mu\nu}, \quad (33)$$

and

$$c_{ZZ}^{\mathcal{F}} = \sum_{\mathcal{F}} \frac{\alpha N_c}{3\pi s_W^2} (g_L^2 + g_R^2) \frac{M_{H_3}}{M_{\mathcal{F}}} y_{\mathcal{F}}. \quad (34)$$

The scalar loop contribution is

$$-\frac{1}{4} \left[1 + \frac{\lambda'^2}{48\pi^2} \log \frac{M_S^2(\mathcal{H})}{\mu^2} + \frac{\eta^2}{192\pi^2} \log \frac{M_V^2(\mathcal{H})}{\mu^2} \right] Z_{\mu\nu} Z^{\mu\nu}, \quad (35)$$

and

$$c_{ZZ}^S = - \sum_S \left(\frac{\lambda'^2}{24\pi^2} + \frac{\eta^2}{96\pi^2} \right) \frac{M_{H_3}}{M_S} y_S. \quad (36)$$

The decay of $H_3 \rightarrow W^+ W^-$ can be easily obtained from the result of $H_3 \rightarrow ZZ$ decay. We consider the couplings between W -boson and new particles are same as Eq. 30, then

$$c_{WW}^V = c_{ZZ}^V, \quad c_{WW}^F = c_{ZZ}^F, \quad c_{WW}^S = c_{ZZ}^S. \quad (37)$$

D. Loop induced $H_3 \rightarrow Z\gamma$

The decay of $H_3 \rightarrow Z\gamma$ can also occur through a set of heavy fermions, scalars and vector bosons. The contribution of the vector boson loop is

$$-\frac{1}{2} \left[1 - \frac{7e^2 c_W}{16\pi^2 s_W} \kappa \kappa' \log \frac{M_V^2(\mathcal{H})}{\mu^2} \right] F_{\mu\nu} Z^{\mu\nu}, \quad (38)$$

the fermion loop contribution is

$$-\frac{1}{2} \left[1 + \frac{e^2 N_c Q_F}{24\pi^2 s_W} (g_L + g_R) \log \frac{M_F^2(\mathcal{H})}{\mu^2} \right] F_{\mu\nu} Z^{\mu\nu}, \quad (39)$$

and the scalar loop contribution is

$$-\frac{1}{2} \left[1 - \frac{\lambda \lambda'}{48\pi^2} \log \frac{M_S(\mathcal{H})}{\mu^2} \right] F_{\mu\nu} Z^{\mu\nu}. \quad (40)$$

We thus obtain

$$\begin{aligned} c_{Z\gamma}^V &= \sum_V \frac{7\alpha c_W}{2\pi s_W} \kappa \kappa' \frac{M_{H_3}}{M_V} y_V, \\ c_{Z\gamma}^F &= - \sum_F \frac{\alpha N_c Q_F}{3\pi s_W} (g_L + g_R) \frac{M_{H_3}}{M_F} y_F, \\ c_{Z\gamma}^S &= \sum_S \frac{\lambda \lambda'}{24\pi^2} \frac{M_{H_3}}{M_S} y_S. \end{aligned} \quad (41)$$

E. The 750 GeV Diphoton excess

Equipped with all the above analytical results, we calculate all the partial widths of the H_3 scalar and the cross section of $pp \rightarrow H_3 \rightarrow \gamma\gamma$ using Eq. 17. We choose $M_{H_W} = M_{H_V} = M_{H_Y} = 800$ GeV as our benchmark parameters. Figure 2 displays the cross section contour of $pp \rightarrow H_3 \rightarrow \gamma\gamma$ at the 13 TeV LHC in the plane of λ_{eff} and v_3 , where λ_{eff} in the limit of $v_3 \gg v_{1,2}$ is given by

- $H_3 H_W H_W : 2kv_2/\sqrt{v_1^2 + v_2^2} + \lambda_{13}v_2^2/(v_1^2 + v_2^2) + \lambda_{23}$
- $H_3 H_Y H_Y : \lambda_{23} + \lambda'_{23}$
- $H_3 H_V H_V : \lambda_{13} + \lambda'_{13}$

For simplicity, we use a universal effective coupling λ_{eff} to represents the above three couplings. The shaded region represents the parameter space allowed by the combined constraint

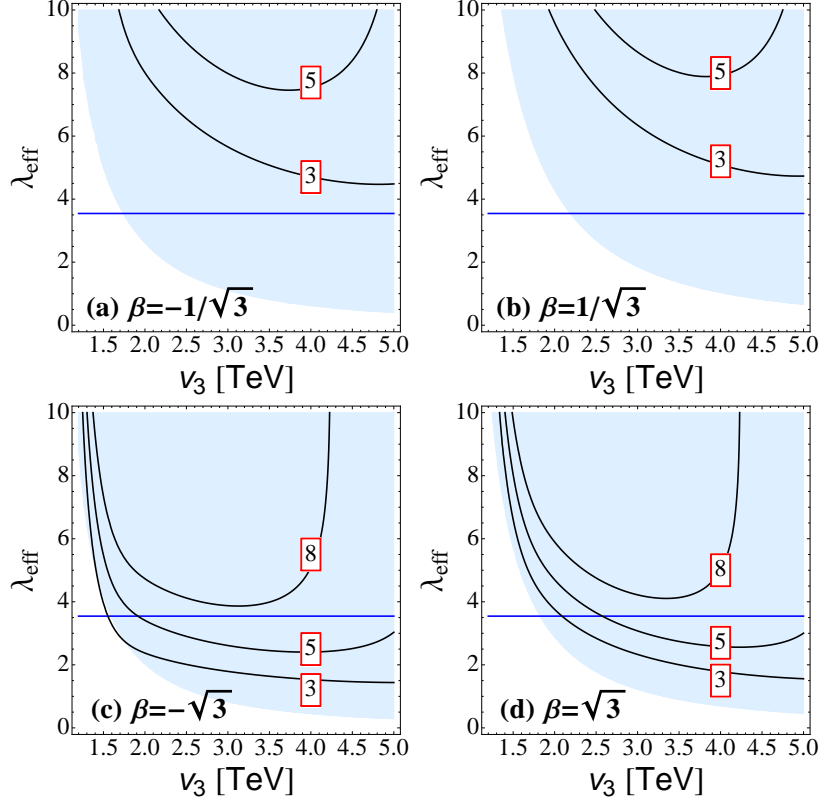


FIG. 2. The contour of $\sigma(pp \rightarrow H_3 \rightarrow \gamma\gamma)$ (fb) in the plane of λ_{eff} and v_3 . The shaded region is consistent with the combined bound of the $WW, ZZ, Z\gamma$, and gg productions at the LHC Run-I while the blue horizontal line represents the perturbation upper limit.

from the $WW, ZZ, Z\gamma$ and gg productions at the LHC Run-1 [120]. The di-jet constraint can be satisfied over the entire parameter space shown in Fig. 2. The blue horizontal line denotes the perturbative bound $\lambda_{\text{eff}} \leq \sqrt{4\pi}$. The cross section for $\beta = \pm\sqrt{3}$ is larger than that for $\beta = \pm 1/\sqrt{3}$ due to the electric charge enhancement of the charged scalars; see Table I. For the models of $\beta = \pm 1/\sqrt{3}$, one cannot obtain enough rate of the diphoton rate to explain the excess even with a large λ_{eff} close to the perturbative limit. For a small v_3 , say $v_3 \leq 1150$ GeV, $M_V = M_Y \leq M_{H_3}/2$. In such a case the decay channels of $H_3 \rightarrow VV/YY$ open at the tree level. That increases the total width of the H_3 scalar and reduces the branching ratio of $H_3 \rightarrow \gamma\gamma$ significantly. A fairly large λ_{eff} is needed to overcome the suppression. Thus, a large v_3 is preferred such that the VV and YY modes are forbidden by kinematics. However, for a large v_3 , the heavy quarks and leptons *alone* cannot produce enough production rate of diphoton pairs to explain the excess; for example, $\sigma(pp \rightarrow H_3 \rightarrow \gamma\gamma) \sim 1$ fb for $v_3 \sim 2$ TeV and $\beta = \pm\sqrt{3}$. We emphasize that one has to include the loop contributions of charged scalars to obtain a sizable rate of diphoton productions. Figure 2(c) and (d) show that the H_3 scalar in the models of $\beta = \pm\sqrt{3}$ could explain the diphoton excess at a large range of parameter space. Unfortunately, the other two models of $\beta = \pm 1/\sqrt{3}$ cannot produce enough diphoton pairs in the entire parameter space; see Figs. 2(a) and (b).

Finally, we comment on the 2 TeV diboson excess observed at the LHC Run-I. The LHC Run-II data do not confirm the excess, but the data do not significantly contradict it either. One has to await for more data to make affirmative conclusion. The Z' boson is predominantly produced via the quark annihilation. It then could decay into a pair of W bosons through the mixing with the Z boson [121]. Figure 3 displays the cross section of $\sigma(pp \rightarrow Z' \rightarrow W^+W^-)$ at the 8 TeV as a function of v_3 or $M_{Z'} (\simeq 0.97v_3)$ in the model of $\beta = -\sqrt{3}$. For a 2 TeV Z' , the cross section of the diboson production is about 3 fb which is consistent with the LHC Run-I data.

IV. CONCLUSION

In this work we explained both the 750 GeV diphoton and the 2 TeV diboson anomalies in the so-called 331 model which has the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge symmetry. The model consists of necessary new ingredients to explain both anomalies, e.g. new heavy

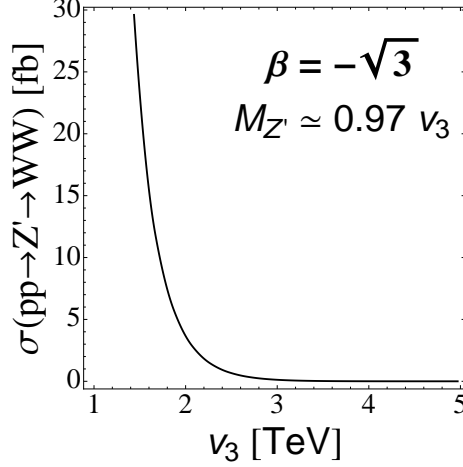


FIG. 3. The cross section of $pp \rightarrow Z' \rightarrow W^+W^-$ as a function of v_3 in the model of $\beta = -\sqrt{3}$ at the 8 TeV LHC.

scalars, heavy quarks and leptons, and new gauge bosons. The symmetry breaking is induced by three scalar triplets. One of them, the scalar triplet χ , is responsible for breaking the $SU(3)_L \otimes U(1)_X$ gauge symmetry down to the SM electroweak symmetry $SU(2)_L \otimes U(1)_Y$ at the scale v_3 . The other two scalar triplets, ρ and η , are in charge of the symmetry breaking of $SU(2)_L \otimes U(1)_Y \rightarrow U(1)_{\text{em}}$ at the electroweak scale. The scalar H_3 , arising from the χ scalar triplet, couples weakly to the SM particles at the tree level. It however can couple to the SM gluon and photons via the loops involving heavy quarks and leptons, charged scalars, and charged vector bosons. We use low energy theorem to calculate effective coupling of $H_3 gg$, $H_3 \gamma\gamma$, $H_3 ZZ$, $H_3 WW$ and $H_3 Z\gamma$. The analytical results can be applied to new physics models satisfying the low energy theorem. We show that H_3 is a good candidate of the 750 GeV diphoton resonance. We also demonstrate that the 2 TeV Z' boson can explain the diboson excess observed at the LHC Run-I. It would be interesting to see if future LHC data will confirm or rule out the two anomalies.

ACKNOWLEDGMENTS

The work is supported in part by the National Science Foundation of China under Grand No. 11275009.

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